

ON THE NATURE OF PHOTOSPHERIC OSCILLATIONS IN STRONG X–RAY BURSTS

Iosif Lapidus ^{1,3}, Luciano Nobili ² and Roberto Turolla ²

¹ Institute of Astronomy, University of Cambridge

Madingley Road, Cambridge CB3 0HA, UK

² Department of Physics, University of Padova

Via Marzolo 8, 35131 Padova, Italy

³ The Royal Astronomical Society Sir Norman Lockyer Fellow

ABSTRACT

A possible sound origin for the photospheric oscillations in the X–ray bursting sources 1608-522 and 2127+119 is suggested. It is shown that standing sound waves in an expanding spherical envelope can have periods very close to the observed ones. The quite large ratio, ~ 10 , of the periods in the two sources is explained in terms of different wave regimes. The relevance of sound oscillations to the observed QPO in type II bursts of the Rapid Burster is also discussed.

Subject headings: stars: individual (1608-522, 2127+119, the Rapid Burster 1730-335) – stars: neutron – stars: oscillations – X–rays: bursts

Accepted for publication in ApJ Letters

1. INTRODUCTION

Photospheric expansion is commonly observed during very powerful X-ray bursts and is widely believed to be associated with a supersonic outflow, driven by the super-Eddington flux produced in the thermonuclear He-burning at the base of the envelope (see e.g. Lewin, Van Paradijs & Taam 1993 for a review on both theoretical and observational aspects of X-ray bursters). In a very recent paper (Nobili, Turolla & Lapidus 1994, hereafter paper I), we presented a more sophisticated model for stationary, radiatively driven winds from X-ray bursting neutron star in which both nuclear burnings at the star surface and Compton heating-cooling in the outflow were self-consistently treated. We proposed that the quasi-stationary phase, during which $L \sim L_{Edd}$, may be regarded as a sequence of steady wind solutions with decreasing \dot{M} which ends when the minimum permitted value of the mass loss rate for the existence of a stationary, supersonic wind is reached. The comparison of our sequence of models with observational data allows for an estimate of both the hardening factor and the mass accretion rate onto the neutron star. This approach has been applied to all sources for which data were available and gave values of $\dot{M}_{acc} \gtrsim 10^{-9} M_{\odot}/yr$, in agreement with expectations (Lapidus, Nobili & Turolla 1994, hereafter paper II).

Additional information on the physics of bursters' envelopes during the expansion phase can be extracted from the second-scale photospheric oscillations registered so far in two sources. In both cases the light curves showed a quite long, \sim tens of seconds, flat top, during which the flux was nearly constant at its maximum level (assumed to coincide with the Eddington limit); oscillations were seen during this phase only. Hakucho observations of 1608-522 have detected a 0.65 s oscillation during a ~ 12 s flat top (Murakami *et al.* 1987) and Ginga data of 2127+119 show a series of oscillations with a

characteristic time scale ~ 7 s during a photospheric expansion phase lasting ~ 30 s (Dotani *et al.* 1990b; Van Paradijs *et al.* 1990). In one of earlier theoretical analysis of such phenomena, the evolution of a gas cloud impinging on the neutron star was studied (Starrfield *et al.* 1982). The infall of the cloud resulted in a sharp burst lasting ~ 1 s, followed by a phase of oscillations with a period of 0.2 s. Even though this period is close to that observed in 1608-522, all the other characteristics of the event were not consistent with observations. A time-dependent study of a largely simplified model of radiatively driven winds (Yahel *et al.* 1984) predicted a typical period for photospheric oscillations of a few tens of milliseconds, and similar results were obtained by McDermott & Taam (1987), who examined the non-radial g-modes in bursting neutron stars atmospheres. A quite sophisticated theory of oscillations was proposed, more recently, by Shibasaki & Ebisuzaki (1989). Although their model successfully reproduced the 0.65 s oscillation seen in 1608-522, it is based on a newtonian hydrostatic picture, while it is widely accepted that a supersonic envelope expansion occurs in very strong bursts.

In this letter we propose that the mechanism responsible for the second-scale oscillations observed in 1608-522 and 2127+119 is the propagation of sound waves. The periods of such waves, computed using the expanding envelope models of paper I, give a correct estimate of the oscillations timescales in both sources. The present model allows also for a simple physical explanation for the different, \sim factor of ten, values of the two observed periods. The relevance of photospheric sound oscillations in connection with QPO in type II bursts from the Rapid Burster is also discussed.

2. SOUND OSCILLATIONS IN TYPE I X-RAY BURSTS

Generally speaking, two main types of sound waves exist in a spherically symmetric medium: radially running waves and standing waves (see e.g. Landau & Lifshitz 1987); the latter form when there is a boundary that can reflect the initial, outward directed wave. In the expanding neutron star envelope, such a role can be played both by the sonic surface and by the photosphere, although the physics is different in the two cases. Below the sonic radius R_s the flow is subsonic and a sound wave can propagate both upstream and downstream, while only a strongly attenuated part of the wave survives at $R > R_s$. In this sense, the initial outgoing disturbance will be mostly reflected backwards near R_s , giving rise to a standing sound wave. What makes the photosphere a peculiar surface, as far as wave propagation is concerned, is the fact that the sound speed drops drastically across it. Below the photosphere, in fact, LTE holds and matter and radiation are strongly coupled, forming a single fluid with pressure $P = P_{gas} + P_{rad} = P_{gas}(1 + \alpha)$. On the contrary, outside the photosphere the radiation is no longer in equilibrium with the gas, no radiative contribution to the total pressure is present, although photons can still exchange momentum with matter, and $P = P_{gas}$. The transition is gradual and it occurs around the radius at which the effective optical depth is about unity. In our model envelopes, the width of the transition layer is $\Delta R \sim R$, so that the approximation of a discontinuous transition may be used only to provide order of magnitude estimates. Clearly temperature and gas density are continuous across the surface, while, because of the change in the equation of state, pressure, energy density and sound speed are not. In particular, the ratio of sound speeds below (region 1) and above (region 2) the photosphere is

$$\frac{v_s^{(1)}}{v_s^{(2)}} = \left[1 + \frac{32\alpha^2}{5(1+8\alpha)} \right]^{1/2},$$

which implies a reflection coefficient for acoustic waves (see again Landau & Lifshitz)

$$\mathbf{R} \sim \left(\frac{v_s^{(1)}/v_s^{(2)} - 1}{v_s^{(1)}/v_s^{(2)} + 1} \right)^2. \quad (1)$$

Here we denote by R_{ph} the lower border of the transition layer, i.e. the radius at which LTE begins to break, and where the ratio of the radiation energy density to the blackbody one, u/aT^4 , starts to drop below unity. The sound speed at R_{ph} is $v_s(R_{ph}) = v_s^{(1)} \propto \sqrt{\alpha T}$ (in all our models $\alpha(R_{ph}) \gg 1$) while, at larger radii, $v_s(R) = v_s^{(2)} \propto \sqrt{T}$.

Before examining in some more details the different wave propagation modes, we present some parameters of the wind models, calculated in paper I, which are relevant to our present analysis. Results are summarized in table 1 where, for each \dot{M} , the following quantities are listed: the sonic radius, R_s , the “sonic” sound crossing time $R_s/v_s(R_s)$, the photospheric radius R_{ph} , computed assuming that $u/aT^4 = 0.9$, $\alpha(R_{ph})$, and the corresponding “photospheric” sound crossing time, $R_{ph}/v_s(R_{ph})$; all the data refer to a neutron star of mass $M_* = 1.5 M_\odot$, radius $R_* = 13.5$ km and to an envelope with nearly solar chemical composition. As can be seen from the table, the photosphere is always inside the sonic radius for wind models with $\dot{M} < 50 \dot{M}_{Edd}$ and, in both 1608-522 and 2127+119, the mass loss rate at the beginning of the quasi-stationary wind phase was estimated to be below this value (see paper II). We mention that the photospheric radii as defined here exceed slightly ($\lesssim 10\%$) those of paper I, where the photosphere was taken to be at $\tau_{eff} = 3$.

Let us consider now the possible regimes of sound oscillations. For a standing wave to be formed, the initial spherical diverging wave must reach the photosphere, reflect back and return to the base of the envelope, so it will take, at least, a time $t_{ph} \sim 2R_{ph}/v_s^{(1)}$ for this standing wave to get stabilized. Data in table 1 show that t_{ph} is shorter than the flat-top durations and “photospheric” standing waves can settle in both sources. The possible occurrence of the second type of standing waves, the “sonic” one, is restricted to 2127+119, because only in this source $t_s \sim 2R_s/v_s(R_s)$ is shorter than the flat-top duration. In order to derive a simple estimate of the pulsational eigenfrequencies, we assume that the local sound speed is constant over the entire shell $R_{in} < R < R_{out}$ and that R_{in} can be taken to be zero to all practical purposes. Here $R_{in} = R_*$, $R_{out} = R_{ph}$ (for “photospheric” waves) and $R_{out} = R_s$ (for “sonic” waves). Both these hypothesis are reasonable since numerical wind solutions show that v_s does not vary by orders of magnitude and $R_* \ll R_{out}$ (see also Table 1). The velocity potential for a monochromatic standing spherical wave has the form (Landau & Lifshitz)

$$\varphi = A \frac{\sin kr}{r} e^{-i\omega t} \quad (2)$$

and from this expression the eigenfrequencies equation can be easily obtained once a boundary condition at $r = R_{out}$ is specified. The fixed boundary condition, $v = \partial\varphi/\partial r = 0$ at $r = R_{out}$, results in

$$\tan z = z, \quad (3)$$

where $z = \omega R_{out}/v_s$. The corresponding periods are

$$P_n^{fix} \sim \frac{4}{1+2n} \cdot \frac{R_{out}}{v_s}, \quad n = 1, 2, \dots \quad (4)$$

In the case of a free oscillating boundary, $\delta p = -\varrho \partial \varphi / \partial t = 0$ at $r = R_{out}$, the eigenfrequencies equation becomes

$$\sin z = 0, \quad (5)$$

and the periods of the eigenmodes are

$$P_n^{free} = \frac{2}{n} \cdot \frac{R_{out}}{v_s}, \quad n = 1, 2, \dots \quad (6)$$

For the burst in 1608-522 observed with Hakucho the mass loss rate at the beginning of the wind phase, i.e. when the luminosity has just reached the Eddington level, was estimated to be $\sim 26 \dot{M}_{Edd}$ (paper II). From the table it can be seen that the sound crossing time ranges from ~ 0.4 to ~ 0.1 s, as the mass of the envelope decreases during the wind phase. The free boundary condition at the photosphere seems to be more appropriate than the fixed one, since there is no physical mechanism to keep the photosphere fixed in space. The period of the principal mode is $P_1 \sim 2R_{ph}/v_s(R_{ph}) \sim 0.8 - 0.1$ s, which is a fairly good estimate for the ~ 0.65 s oscillations observed in 1608-522. Although not only the principal mode may be present, the amplitudes of higher order oscillations are much more reduced by damping: in fact, denoting by γ_n the n -th mode sound absorption coefficient, it is $\gamma_n/\gamma_1 = \omega_n^2/\omega_1^2 = n^2$.

In 2127+119 photospheric oscillations with a characteristic time scale ~ 7 s were observed with Ginga during a ~ 30 s flat top of the light curve. In paper II, the initial envelope mass was estimated to be $M_{env} \sim 8 \times 10^{22}$ g, corresponding to $\dot{M} \sim 16 \dot{M}_{Edd}$. Now the flat top duration is longer than t_s , so it is reasonable to expect that also a “sonic” standing wave has time to develop because, for a typical value of $\alpha \sim 800$, about 15 % of the initial disturbance is transmitted through the photosphere (see eq. [1]) and will form later the “sonic” standing wave. We note that now both the free and the fixed boundary

conditions provide radial oscillations of the photosphere, since, for the fixed one, the node is at R_s and not at R_{ph} . Equations (4) and (6) show, however, that the differences in the eigenfrequencies are relatively small, so, despite our ignorance about the exact form of the boundary condition at R_s , we can get a fair estimate of the period. Referring to the free-boundary spectrum, the fundamental mode has $P_1 \sim 22$ s and can hardly be detected, being too close to the total duration of the flat-top phase. The second harmonic has a period $P_2 \sim 11$ s which is close to the observed ~ 7 s periodicity; of course this mode is weaker than the fundamental one, but it should dominate the Fourier spectrum in the absence of the main one. Although the simultaneous detection of both kinds of standing waves in 2127+119 cannot be excluded, short wavelength “photospheric” modes are more severely damped with respect to the longer “sonic” modes.

Another observational fact which supports the hypothesis of radial photospheric oscillations is the phase lag of about 180° between the oscillations in the low- and high-energy bands. The constant luminosity ($L \simeq L_{Edd}$) radial oscillations of the photosphere result in simultaneous oscillations of the photospheric temperature and the degree of Comptonization of radiation crossing the translucent region, and therefore in the simultaneous oscillations of the whole spectrum registered by a distant observer. As the temperature reaches its maximum, the counts in the hard bands are the maximum, while those in the soft bands are the minimum, and vice versa. It means that the strong anticorrelation between the counts at low and high energies oscillating with the same frequency, provides the additional evidence in support of our theory of sound waves.

3. QPO IN TYPE II X-RAY BURSTS

The physics of “abnormal” type II bursts is still unclear, but the currently dominating idea is that these bursts are due to some sort of instability in the inner regions of the accretion flow onto the neutron star. A detailed review of the observations of the Rapid Buster (RB) can be found in Lewin *et al.* Type II bursts, too, show some observational evidence of photospheric expansion. Kunieda *et al.* (1984) and Tan *et al.* (1991), in fact, have shown that there is a correlation between the color temperature, T_{col} , and the luminosity, L , observed during the flat tops of relatively long bursts. It was found that $L \propto T_{col}^6$ and this led Tan *et al.* to suggest that the radiation released during a type II burst comes from a photosphere whose radius is larger for higher accretion rates. Quasi-periodic oscillations (QPO) with frequencies ~ 2 Hz were discovered in type II bursts from the RB with Hakicho, and a further analysis of these QPO indicated that, most probably, they have no relation to other forms of QPO observed in many bright LMXBs (see e.g. Lewin *et al.* for references). An indirect confirmation of this fact was provided by Rutledge *et al.* (1993), who have shown that the RB is neither an atoll nor a Z source in the classification of Hasinger & Van der Klis (1989). Lubin *et al.* (1991) showed that the QPO centroid frequencies in type II bursts range from ~ 2 to ~ 7 Hz and they are strongly anticorrelated with the peak flux in the burst. It was proposed by Dotani *et al.* (1990a) that QPO can be described by changes in the photospheric radius and temperature. Lewin *et al.* (1991) have given some arguments why this possibility is more likely than the changes in the temperature of a blackbody emitter with constant apparent area.

Here we suggest that QPO in type II bursts could have the same origin as photospheric oscillations in “normal”, type I bursts. We emphasize, however, that no quasi-stationary, Eddington luminosity phase is attained in type II

bursts, so our wind models are not directly applicable for describing the expansion/contraction phase of the Rapid Burster. Still, it would be reasonable to assume that, if the luminosity is sub–Eddington or super–Eddington during a short unstationary stage, the structure of the outflow is similar to that of our models representing the final part of the expansion/contraction phase, i.e. those ones in the lower end of the M_{env} (or \dot{M}) range. The characteristics of these models may provide order of magnitude estimates for various quantities in the expanded RB envelope. The frequency of the principal mode of the “photospheric” standing wave is ≈ 10 Hz for the models with the lowest possible \dot{M} (last lines in Table 1), and ≈ 2 Hz for an envelope with $\dot{M} \sim 20\dot{M}_E$, as that one of both 1608-522 and 2127+119 at the beginning of the expansion/contraction phase. We see that the ~ 2 -7 Hz range of the QPO centroid frequencies of the RB may be easily reproduced in frame of our hypothesis. Let us consider now how the sound frequency correlates with the peak flux in the burst. In our wind solutions the larger is the envelope mass, the larger is the photospheric radius and the lower is the temperature at the photosphere. As it may be seen from the last column of Table 1, the period of the “photospheric” standing waves monotonically decreases with decreasing \dot{M} . The sound oscillations frequency, therefore, will be higher for the less energetic bursts and this is indeed consistent with the observed anticorrelation of the QPO frequency with the average burst peak flux. It is also not surprising that, being of sound origin, the QPO in the Rapid Burster are qualitatively different from the standard QPO seen in many bright LMXBs which are probably due to dynamical effects.

4. CONCLUSIONS

We have presented a simple model for the photospheric oscillations observed during type I X-ray bursts from 1608-522 and 2127+119, and in type II bursts of the Rapid Burster. In particular, we have shown that the periods of standing sound waves in the expanding envelopes analyzed in paper I are very close to the observed ones. The quite large, about a factor of 10, difference in the observed periods is explained in terms of the different wave regimes in the two sources: “photospheric” and “sonic” standing sound waves. The acoustic wave hypothesis could also explain the QPO in type II bursts from the Rapid Burster; both the QPO centroid frequencies and the strong anticorrelation of the centroid frequency with the peak flux in the burst are, in fact, consistent with the “photospheric” standing sound wave scenario.

The existence of two different wave regimes can lead to significant observational predictions. A number of bursters exhibits, in fact, strong bursts with precursors which have long expansion phases. The “sonic” standing wave has, therefore, enough time to form and oscillations with periods up to ~ 25 s may be revealed in the light curves flat tops. Both the re-examination of archive data and new observations with X-ray telescopes on board satellites of the current generation, like ASCA, may definitely access the existence of such oscillations.

ACKNOWLEDGMENTS

We are indebted to an anonymous referee whose constructive criticism helped in improving an earlier version of this paper.

REFERENCES

Dotani, T., Mitsuda, K., Inoue, H., Tanaka, Y., Kawai, N., Tawara, Y., Makishima, K., Van Paradijs, J., Penninx, W., Van der Klis, M., Tan, J., & Lewin, W.H.G. 1990a, *ApJ*, 350, 395

Dotani, T., Inoue, H., Murakami, T., Nagase, F., Tanaka, Y., Tsuru, T., Makishima, K., Ohashi, T., & Corbet, R.H.D. 1990b, *Nature*, 347, 534

Hasinger, G., & Van der Klis, M. 1989, *A&A*, 225, 79

Kunieda, H., Tawara, Y., Hayakawa, S., Masai, K., Nagase, F., Inoue, H., Koyama, K., Makino, F., Makishima, K., Matsuoka, M., Murakami, T., Oda, M., Ogawara, Y., Ohashi, T., Shibasaki, N., Tanaka, Y., Kondo, I., Miyamoto, S., Tsunemi, H., & Yamashita, K. 1984, *PASJ*, 36, 215

Landau, L.D., & Lifshitz, E.M. 1987, *Fluid Mechanics*, (Oxford: Pergamon)

Lapidus, I., Nobili, L., & Turolla, R. 1994, *ApJ*, in press (Paper II)

Lewin, W.H.G., Lubin, L.M., Van Paradijs, J., & Van der Klis, M. 1991, *A&A*, 248, 538

Lewin, W.H.G., Van Paradijs, J., & Taam, R.E. 1993, *Space Sci. Rev.*, 62, 223

Lubin, L.M., Stella, L., Lewin, W.H.G., Tan, J., Van Paradijs, J., Van der Klis, M., & Penninx, W. 1991, *MNRAS*, 249, 300

McDermott, P.N., & Taam, R.E. 1987, *ApJ*, 318, 278

Murakami, T., Inoue, H., Makishima, K., & Hoshi, R. 1987, *PASJ*, 39, 879

Nobili, L., Turolla, R., & Lapidus, I. 1994, *ApJ*, in press (Paper I)

Rutledge, R.E., Lubin, L.M., Lewin, W.H.G., Van Paradijs, J., & Van der Klis, M. 1993, *MNRAS*, submitted

Shibasaki, N., & Ebisuzaki, T. 1989, *PASJ*, 41, 641

Starrfield, S., Kenyon, S., Sparks, W.M., & Truran, J.W. 1982, *ApJ*, 258, 683

Tan, J., Lewin, W.H.G., Lubin, L.M., Van Paradijs, J., Penninx, W., Damen, E., & Stella, L. 1991, *MNRAS*, 251, 1

Van Paradijs, J., Dotani, T., Tanaka, Y., & Tsuru, T. 1990, PASJ, 42, 633

Yahel, R.Z., Brinkmann, W., & Braun, A. 1984, A&A, 139, 359

Table 1
Parameters of Model Atmospheres

\dot{M} (\dot{M}_{Edd})	R_s (10^3 km)	$R_s/v_s(R_s)$ (s)	R_{ph} (10^3 km)	$\alpha(R_{ph})$ (10^2)	$R_{ph}/v_s(R_{ph})$ (s)
50.2	5.9	27.9	2.20	4.6	1.47
26.7	4.0	14.9	0.96	6.6	0.39
19.8	3.3	10.9	0.65	7.8	0.21
12.7	3.9	11.4	0.49	10.3	0.10
8.8	5.0	13.4	0.27	11.2	0.06
6.7	6.0	15.2	0.23	13.3	0.05
5.9	7.0	17.3	0.21	13.8	0.04
4.8	9.5	23.5	0.19	13.9	0.04
2.8	16.4	38.0	0.16	14.8	0.03